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Splitting the pomeron into two jets: a novel process at HERA

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Abstract

We study a novel property of large rapidity-gap events in deep inelastic scattering at HERA: splitting the pomeron into two jets in the photon-pomeron fusion reaction $\gamma^* \mathbf{I\!P} \rightarrow q\bar{q}$. It gives rise to the diffraction dissociation of virtual photons into the back-to-back jets. We find that at large invariant mass M of two jets, $M^2 \gg Q^2$, the transverse momentum of jets comes from the intrinsic transverse momentum of gluons in the pomeron, and the photon-pomeron fusion direcly probes the diffrential gluon structure function of the proton $\partial G(x,q^2)/\partial \log q^2$ at the virtuality $q^2 \sim k^2$. We present estimates for the jet production cross section, which show the process is easily measurable at HERA.

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Diffraction dissociation reactions $a+p \to X+p'$ can be interpreted as interaction of the projectile b with the pomeron $a+\mathbb{P} \to X$ [1]. When the projectile is the virtual photon γ^*, then at moderately large invariant mass M of the diffractively excited state, $M^2 \sim Q^2$, one probes the valence $q\bar{q}$ component of the pomeron (Q^2 is the virtuality of the photon). The corresponding valence structure function of the pomeron was evaluated in [2] and the predicted rate of the diffraction dissociation is consistent with the first data from the ZEUS experiment at HERA [3].

In [2] we have also predicted diffraction dissociation of photons into back-to-back jets, alias the splitting of pomerons into two jets, the partonic subprocess for which is the (virtual) photon-pomeron fusion

$$\gamma^* + \mathbb{P} \to q + \bar{q}\,. \tag{1}$$

This process has a very clean signature: the (anti)quarks produced with large transverse momentum k with respect to the photon-pomeron collision axis, materialize as the back-to-back jets in the photon-pomeron c.m.s. Understanding to which extent the pomeron can be treated as a particle is one of the pressing issues in the pomeron physics [2,4], and the salient feature of this "two-body" reaction (1) is precisely that here the pomeron participates as an "elementary" particle, which carries 100 per cent of its momentum. On the other hand, the QCD pomeron describes the colour-singlet exchange by (the two) gluons and, as such, it is definitely not an elementary particle. In this paper we update the considerations [2] and discuss in detail how this gluonic substructure of the pomeron manifests itself in the back-to-back jet production. Our major finding is that, in the most interesting kinematical domain, the reaction (1) proceeds via interaction of the photon with the two gluons of the pomeron in such a way that the transverse momenta of the quark and antiquark jets come entirely from the intrinsic transverse momentum of gluons in the pomeron. Consequently, the underlying process can be dubbed the splitting of pomerons into two jets. We demonstrate how the diffraction dissociation of photons into the back-to-back jets direcly probes the differential gluon structure function of the proton $f(x,q^2) = \partial G(x,q^2)/\partial \log q^2$ at the virtuality $q^2 \sim k^2$.

Experimentally, one studies the reaction $\gamma^* + p \to X + p$, which can be viewed as emisssion of the pomeron $\mathbb{P}$ by the proton followed by the photoabsorption $\gamma^* + \mathbb{P} \to X$.

By definition of the diffraction dissociation, the fraction $x_{\bf IP}$ of proton's momentum carried by the pomeron is small: $x_{\bf IP} \lesssim (0.05\text{-}0.1)$. The recoil proton emerges in the final state separated from the debris of the diffractively excited photon by a large (pseudo)rapidity gap

$$\Delta\eta \approx \log\left(\frac{1}{x_{\bf IP}}\right) . \tag{2}$$

The invariant mass M of the diffractively excited state and the total mass squared s of the hadronic final state are related by $M^2 + Q^2 = x_{\bf IP}(s + Q^2)$. Hereafter, for the sake of simplicity, we consider the forward diffraction dissociation, with the vanishing transverse momentum $\vec{p}_\perp$ of the recoil proton: $t = -\vec{p}_\perp^{\,2} = 0$; extension of our results to finite, but small $t \ll k^2$ poses no problems. We consider the jet cross section at the partonic level, i.e., we calculate the differential cross section for production of the high-k (anti)quarks in the photon-pomeron fusion (1). The diffraction dissociation cross section is described by diagrams of Fig.1.

The staring point of our analysis is the generalization of formulas for the corresponding differential cross section $d\sigma_D(\gamma^*_{T,L} \to q\bar{q})/dt$ derived by us in [2,5]. For the diffraction dissociation of the (γ^*_T) transverse and (γ^*_L) longitudinal photons we obtain

$$\begin{aligned}
\left.\frac{d\sigma_D(\gamma^*_T \to q\bar{q})}{dt}\right|_{t=0} &= \frac{\alpha_{em}}{6\pi}\sum e_f^2 \int_0^1 dz \int d^2\vec{k}\; d^2\vec{\kappa}\; d^2\vec{\kappa}'\; \alpha_S^2 \cdot \frac{1}{\kappa^4}\frac{\partial G(x_g,\kappa^2)}{\partial \log\kappa^2}\cdot\frac{1}{\kappa'^4}\frac{\partial G(x_g,\kappa'^2)}{\partial\log\kappa'^2} \\
&\left\{ \frac{[z^2+(1-z)^2]\vec{k}^2 + m_f^2}{[\vec{k}^2+\varepsilon^2]^2} - \frac{[z^2+(1-z)^2]\vec{k}\cdot(\vec{k}+\vec{\kappa}) + m_f^2}{[\vec{k}^2+\varepsilon^2][(\vec{k}+\vec{\kappa})^2+\varepsilon^2]} \right. \\
&\left. -\frac{[z^2+(1-z)^2]\vec{k}\cdot(\vec{k}-\vec{\kappa}') + m_f^2}{[\vec{k}^2+\varepsilon^2][(\vec{k}-\vec{\kappa}')^2+\varepsilon^2]} + \frac{[z^2+(1-z)^2](\vec{k}+\vec{\kappa})\cdot(\vec{k}-\vec{\kappa}') + m_f^2}{[(\vec{k}+\vec{\kappa})^2+\varepsilon^2][(\vec{k}-\vec{\kappa}')^2+\varepsilon^2]} \right\} ,
\end{aligned} \tag{3}$$

$$\begin{aligned}
\left.\frac{d\sigma_D(\gamma^*_L \to q\bar{q})}{dt}\right|_{t=0} &= \frac{\alpha_{em}}{6\pi}\sum e_f^2 \int_0^1 dz\; 4Q^2 z^2(1-z)^2 \int d^2\vec{k}\; d^2\vec{\kappa}\; d^2\vec{\kappa}' \\
\alpha_S^2 \cdot \frac{1}{\kappa^4}\frac{\partial G(x_g,\kappa^2)}{\partial \log\kappa^2}\cdot\frac{1}{\kappa'^4}\frac{\partial G(x_g,\kappa'^2)}{\partial\log\kappa'^2} &\left\{ \frac{1}{[\vec{k}^2+\varepsilon^2]^2} - \frac{1}{[\vec{k}^2+\varepsilon^2][(\vec{k}+\vec{\kappa})^2+\varepsilon^2]} \right. \\
&\left. -\frac{1}{[\vec{k}^2+\varepsilon^2][(\vec{k}-\vec{\kappa}')^2+\varepsilon^2]} + \frac{1}{[(\vec{k}+\vec{\kappa})^2+\varepsilon^2][(\vec{k}-\vec{\kappa}')^2+\varepsilon^2]} \right\} .
\end{aligned} \tag{4}$$

Here e_f is the quark charge in units of the electron charge, m_f is the quark mass, z is a fraction of the photon's lightcone momentum carried by the (anti)quark, and $\vec{k},\vec{\kappa},\vec{\kappa}'$ are the transverse momenta of the produced quark and of gluons in pomerons (Fig.1). The

observed back-to-back jets will have the transverse momenta $\vec{k}$ and $-\vec{k}$. The invariant mass M of the two-jet final state equals

$$M^2 = \frac{m_f^2 + \vec{k}^2}{z(1-z)} \tag{5}$$

and

$$dz\, dk^2 = dM^2\, dk^2 \frac{m_f^2 + k^2}{M^4} \left(1 - 4\frac{m_f + k^2}{M^2}\right)^{-1/2} , \tag{6}$$

where the factor

$$\frac{1}{\cos\theta} = \left(1 - 4\frac{m_f + k^2}{M^2}\right)^{-1/2} \tag{7}$$

corresponds to the familiar Jacobian-peak at the jet production angles $\theta \sim 90^o$ in the photon-pomeron c.m.s. In the Born approximation, the differential gluon density is related to the two-body formfactor of the nucleon $\langle N|\exp(i\vec{\kappa}_1\vec{r}_1 + i\vec{k}_2\vec{r}_2)|N\rangle$ by the equation [6]

$$\frac{1}{\kappa^4}\frac{\partial G(x,\kappa^2)}{\partial \log\kappa^2} = \frac{4\alpha_S(\kappa^2)}{\pi(\kappa^2+\mu_G^2)^2}\left(1 - \langle N|\exp(i\vec{\kappa}(\vec{r}_1 - \vec{r}_2)|N\rangle\right) , \tag{8}$$

and the limiting form of Eqs. (3,4) derived in [2,4] is obtained. (In (8), $\mu_G = R_c^{-1}$ and R_c is the correlation radius for the perturbative gluons [2,4-9]).

After integrations over the azymuthal angles of the gluon momenta $\vec{\kappa}$ and $\vec{\kappa}'$, the differential cross sections can be written in the compact form

$$\left.\frac{d\sigma_D(\gamma_T^* \to q\bar{q})}{dM^2 dk^2 dt}\right|_{t=0} = \frac{\pi^2\alpha_{em}}{6\cos\theta}\cdot\frac{m_f^2+k^2}{M^4}\alpha_S^2(k^2)\left\{\left(1 - 2\frac{k^2+m_f^2}{M^2}\right)\Phi_1^2 + m_f^2\Phi_2^2\right\} , \tag{9}$$

$$\left.\frac{d\sigma_D(\gamma_L^* \to q\bar{q})}{dM^2 dk^2 dt}\right|_{t=0} = \frac{\pi^2\alpha_{em}Q^2}{6\cos\theta}\cdot\frac{(m_f^2+k^2)^3}{M^8}\alpha_S^2(k^2)\Phi_2^2 , \tag{10}$$

where

$$\begin{aligned}\Phi_1 &= \int \frac{d\kappa^2}{\kappa^4} f(x,\kappa^2)\left[\frac{k}{k^2+\varepsilon^2} - \frac{k}{\sqrt{a^2-b^2}} + \frac{2k\kappa^2}{a^2-b^2+a\sqrt{a^2-b^2}}\right] \\ &= \frac{2k\varepsilon^2}{(k^2+\varepsilon^2)^3}\int\frac{d\kappa^2}{\kappa^2}W_1(\frac{k^2}{\varepsilon^2},\frac{\kappa^2}{k^2})f(x,\kappa^2) ,\end{aligned} \tag{11}$$

$$\Phi_2 = \int \frac{d\kappa^2}{\kappa^4} f(x,\kappa^2)\left[\frac{1}{\sqrt{a^2-b^2}} - \frac{1}{k^2+\varepsilon^2}\right] \approx \frac{k^2-\epsilon^2}{(k^2+\varepsilon^2)^3}\cdot\int\frac{d\kappa^2}{\kappa^2}W_2(\frac{k^2}{\varepsilon^2},\frac{\kappa^2}{k^2})f(x,\kappa^2) , \tag{12}$$

$$a = \varepsilon^2 + k^2 + \kappa^2 , \tag{13}$$

$$b = 2k\kappa\,. \tag{14}$$

The useful relations are also

$$\varepsilon^2 = (k^2 + m_f^2)\frac{Q^2}{M^2} + m_f^2\,, \tag{15}$$

$$k^2 + \varepsilon^2 = (k^2 + m_f^2)\frac{M^2 + Q^2}{M^2}\,. \tag{16}$$

The kernels W_i are functions of the dimensionless variables $\omega = k^2/\varepsilon^2$ and $\tau = \kappa^2/k^2$. They have similar properties as a function of ω and τ. Hereafter we concentrate on diffraction dissociation of the transverse photons, because diffraction dissociation of the longitudinal photons into two jets has much smaller cross section [2]. We also can neglect the $m_f^2\Phi_2^2$ term in Eq. (9).

The kernel W_1 is shown in Fig. 2. In deep inelastic scattering, at large Q^2, we have $\omega = M^2/Q^2$. In the excitation of the moderately large masses $M^2 \lesssim Q^2$, the kernel W_1 is essentially flat, has the unity height, and extends up to $\kappa^2 \approx A_T k^2$, where $A_T \approx 3$. Consequently, at $\omega \lesssim 1$,

$$\int d\tau\, W_1(\omega,\tau) f(x,\kappa^2) = G(x, A_T k^2) \tag{17}$$

and

$$\Phi_1 = \frac{2k\varepsilon^2}{(k^2+\varepsilon^2)^3} G(x, A_T k^2)\,, \tag{18}$$

which updates Eq. (60) of [2]. The corresponding differential cross section for the jet production at $M^2 \lesssim Q^2$ equals (here we neglect the last m_f^2 term in ε^2 in the numerator)

$$\left.\frac{d\sigma_D(\gamma_T^* \to q\bar{q})}{dM^2 dk^2 dt}\right|_{t=0} = \sum e_f^2 \frac{\pi^2 \alpha_{em} \alpha_S^2(k^2)}{6\cos\theta}\left(1 - 2\frac{k^2 + m_f^2}{M^2}\right) \tag{19}$$
$$\cdot \frac{1}{M^4}\frac{4k^2}{(k^2+m_f^2)^3}\left(\frac{Q^2}{M^2+Q^2}\right)^2\left(\frac{M^2}{M^2+Q^2}\right)^3 G(x, A_T k^2)^2$$

Because of the logarithmic κ^2 integration, at $\omega \lesssim 1$ the major contribution to the integral for Φ_1 comes from $\kappa^2 \ll k^2$ and we have a semblance of the Leading-LogQ^2 approximation (LLQA). This means that the transverse momentum of gluons of the pomeron contributes little to the transverse momentum of jets. Therefore, in the diffractive excitation to low mass states $M^2 \lesssim Q^2$, the transverse momentum of jets comes predominantly from the

intrinsic transverse momentum of the quark and antiquark in the $q\bar{q}$ Fock state of the photon. The dk^2/k^4 transverse momentum distribution corresponds to a dominance of the diffraction dissociation cross section by the $q\bar{q}$ configurations in which (anti)quarks have small, hadronic, value of the transverse momentum irrespective of the value of Q^2 [2,4].

Much more interesting case is diffraction excitation of heavier masses, $M^2 \gg Q^2$, i.e., $\omega \gg 1$, which requires utilization of the transverse momentum of gluons in the pomeron. Here the kernel develops a very sharp resonance peak at $\tau = \kappa^2/k^2 \sim 1$, which clearly shows that in this regime the transverse momentum of jets comes entirely from the transverse momentum $\vec{\kappa} \approx \vec{k}$ of gluons. The LLQA considerations are completely inapplicable here. This resonance contribution can be quantified as follows: The height of the peak approximately follows the law

$$W_1^{(max)} \sim \frac{\omega}{2}\,, \tag{20}$$

and the width of the peak is approximately constant in the $\log(\kappa^2/k^2)$ scale. A separation of the kernel W_1 into the peak component P_1 and the plateau is somewhat convention dependent. We define the resonance peak distribution via the decomposition

$$W_1(\omega,\tau) = W_1(\omega=1,\tau) + \frac{\omega}{2}P_1(\omega,\tau)\,. \tag{21}$$

The resulting peak distribution function $P_1(\omega,\tau)$ and its variation with ω are shown in Fig. 3. At $\omega >> 1$ it becomes an approximately scaling function of τ. Because $P_1(\omega,\tau)$ is much sharper function of κ^2 than $f(x,\kappa^2)$, we can write

$$\int d\log\tau\, P_1(\omega,\tau) f(x,\kappa^2) = S(\omega) f(x,\tau_T k^2)\,, \tag{22}$$

where $S(\omega)$ is an area under the resonance peak, and τ_T corresponds to the center of gravity of the peak, $\langle\log\tau\rangle = \log\tau_T$. The area under the peak $S(\omega)$ and its center of gravity are shown in Fig. 4. At the asymptotically large $\omega \gg 1$ the resonance area integrates to unity, $S(\omega \gg 1) = 1$, which is a viable first approximation at $\omega \gtrsim 3$. Consequently, at $\omega \gg 1$

$$\Phi_1 = \frac{k}{(k^2+\varepsilon^2)^2}\left[\frac{2G(x,A_T k^2)}{\omega} + S(\omega) f(x,\tau_T k^2)\right] \approx \frac{k}{(k^2+\varepsilon^2)^2} S(\omega) f(x,\tau_T k^2) \tag{23}$$

and

$$\left.\frac{d\sigma_D(\gamma_T^* \to q\bar{q})}{dM^2 dk^2 dt}\right|_{t=0} = \tag{24}$$

$$\sum e_f^2 \frac{\pi^2\alpha_{em}\alpha_S^2(k^2)}{6\cos\theta}\left(1-2\frac{k^2+m_f^2}{M^2}\right)\cdot\frac{1}{M^4}\frac{k^2}{(k^2+m_f^2)^3}\left(\frac{M^2}{M^2+Q^2}\right)^3 f(x,\tau_T k^2)^2\,,$$

which update Eqs. (57,58) of [2]. The emerging possibility of directly measuring the differential gluon distribution is a novel property of the pomeron splitting reaction. For the numerical estimates of $f(x,Q^2)$ in terms of the more familiar integrated gluon distribution function $G(x,Q^2)$, we can use the so-called Double-Leading-Log-Approximation (DLLA) identity, which also holds in the BFKL regime [7] (here $N_f=3$ active flavours are assumed)

$$\frac{3}{4}\cdot\frac{1}{G(\xi,Q^2)}\frac{\partial^2 G(\xi,r)}{\partial\log(1/x)\,\partial\log\log Q^2}=1\,. \tag{25}$$

In [7,8] we have shown that the effective intercept

$$\Delta_{eff}(x,Q^2)=\frac{1}{G(x,Q^2)}\frac{\partial G(x,Q^2)}{\partial\log(1/x)} \tag{26}$$

is a relatively slow function of Q^2 and x, and in the kinematical range of the interest at HERA, $\Delta_{eff}(x,Q^2)\approx\Delta_{\mathbf{IP}}=0.4$ to better than the factor 2. Here $\Delta_{\mathbf{IP}}$ is an intercept of the BFKL pomeron evaluated in [8]. Then the DLLA identity (25) gives an estimate

$$f(x,Q^2)\approx\frac{3\alpha_S(k^2)}{\pi\Delta_{\mathbf{IP}}}G(x,Q^2)\,. \tag{27}$$

This estimate shows that the term $\propto f(x,\xi_T k^2)$ in Eq. (23) dominates already starting with $\omega \gtrsim (2-4)$. Only this term contributes to the real photoproduction, $Q^2=0$, at $k^2\gg m_f^2$ of the interest.

Above we have considered the forward diffraction dissociation $t=0$. At $|t|\ll k^2$, which is the dominant region in diffraction dissociation into two jets, we can write $d\sigma_D/dt = d\sigma_D/dt|_{t=0}\cdot\exp(B_D t)$ with the slope B_D close to the slope of the πN elastic scattering, $B_D\sim 10\mathrm{GeV}^{-2}$. The cross section, integrated over the recoil protons, can be estimated as

$$\int dt\frac{d\sigma_D(\gamma^*\to q+\bar{q})}{dM^2dk^2dt}\approx\frac{1}{B_D}\left.\frac{d\sigma_D(\gamma^*\to q+\bar{q})}{dM^2dk^2dt}\right|_{t=0} \tag{28}$$

In order to have an idea on the magnitude of the pomeron splitting cross section in the real photoproduction, consider the cross section integrated over masses $M^2\gtrsim 4k^2$. Ignoring certain enhancement coming from the Jacobian peak, and putting $\cos\theta\sim 1$, we find

$$\frac{d\sigma_D}{dk^2}\approx\frac{\pi^2\alpha_{em}\alpha_S^2(k^2)}{24B_D}\frac{1}{(k^2+m_f^2)^3}f(x,\tau_T k^2)^2\,, \tag{29}$$

which is dominated by $M^2 \sim 4k^2$. The total cross section for producing jets with the transverse momentum above certain threshold $k^2 > k_0^2$ equals

$$\sigma_D(k^2 > k_0^2) = \int_{k_0^2} dk^2 \frac{d\sigma_D}{dk^2} \approx \frac{\pi^2 \alpha_{em} \alpha_S^2(k_0^2)}{48 B_D} \frac{1}{(k_0^2 + m_f^2)^2} f(x, \tau_T k_0^2)^2 \,. \tag{30}$$

Take $k_0^2 = 5\mathrm{GeV}^2$ in the real photoproduction at the typical HERA center of mass energy squared $s = 4\nu E_p \sim 5 \cdot 10^4 \mathrm{GeV}^2$. The gluon distribution must be evaluated at $x = x_0 \sim M^2/s \sim 4k_0^2/s \approx 5 \cdot 10^{-4}$. Our evaluation [9] based on the BFKL evolution analysis, gives $G(x_0, \tau_T k_0^2) \sim 20$ and Eq. (27) gives $f(x_0, \tau_T k_0^2) \sim 20$. Then $\sigma_D(k^2 \gtrsim 5\mathrm{GeV}^2) \sim 0.1\mu$b. This must be compared with our prediction [2] that the diffraction dissociation of real photons into the $q\bar{q}$ states makes about 15% of the total cross section, i.e., $\sigma_D(tot) \sim 20\mu$b, which is in perfect agreement with the preliminary data from HERA [10].

We conclude that diffraction dissociation of photons into back-to-back jets has sufficiently large cross section to be observable at HERA, and is a potentially important tool for measuring the differential gluon distribution $\partial G(x, Q^2)/\partial \log Q^2$.

A fresh look at the difference of mechanisms of production of large transverse momentum jets at $M^2 \lesssim Q^2$ and $M^2 \gg Q^2$ is in order. Usually, one thinks of production of jets in terms of the elementary partonic subprocesses like the photon-gluon fusion $\gamma^* + g \rightarrow q + \bar{q}$ with production of the large-k quarks. We have shown that in the regime of $M^2 \lesssim Q^2$, gluons in the pomeron have small transverse momenta $\kappa^2 \ll k^2$ and, with certain reservations, here pomeron acts like an elementary particle whose substructure is not important. In the opposite to that, at $M^2 \gg Q^2$, the transverse momenta of the back-to-back jets rather come from the back-to-back transverse momentum of gluons in the pomeron, and the relevant subprocess for splitting the pomeron into two jets is the manifestly three-body, photon-two-gluon collision in the initial state. Here the gluonic substructure of the pomeron enters in a most crucial manner, which precludes any consideration of the pomeron as an elementary particle. Furthermore, it would be completely illegitimate to interpret such a jet production as interaction of the photon whith one parton of the pomeron which carries 100% of the pomeron's momentum. Notice, that in the regime of $M^2 \lesssim Q^2$ the cross section is $\propto G(x, k^2)^2$, whereas at $M^2 \gg Q^2$ it is $\propto [\partial G(x, k^2)/\partial \log k^2]^2$. This striking difference between the two cases is just one of manifestations of the lack of factorization for the QCD

pomeron contribution to the diffraction dissociation [2].

Conclusions:

We found a novel mechanism of diffraction dissociation of photons into two back-to-back jets, in which the transverse momentum of jets originates from the transverse momentum of gluons in the pomeron. This mechanism can not be interpreted in familar terms of hard interaction of the photon with one energetic parton of the pomeron. The corresponding jet production cross section Eq. (25) emerges as a unique direct probe of the differential gluon structure function of the proton $\partial G(x,Q^2)/\partial \log Q^2$. The jet production cross section is large and must easily be measurable at HERA.

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Figure captions:

Fig.1 - One of the 16 perturbative QCD diagrams for the diffraction dissociation cross section. The vertical dashed line shows the unitarity cut corresponding to the diffractively produced state.

Fig.2 - The kernel $W_1(\omega,\tau)$ for different values of ω. The resonance peak at $\tau \sim 0$ developes at large ω.

fig.3 - The scaling properties of the resonanace peak $P_1(\omega,\tau)$.

Fig.4 - The area of the resonance peak $S(\omega)$ and its center of gravity τ_T as a function of ω.